

New near-threshold mesons

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We show that under a number of rather plausible assumptions QCD spectrum may contain a number of mesons which have not been predicted or observed. Such states will have the quantum numbers of two existing mesons and masses very close to the dissociation threshold into the two mesons. Moreover, at least one of the two mesonic constituents itself must be very close to its dissociation threshold. In particular, one might expect the existence of loosely bound systems of D and $D_{sJ}^*(2317)$; similarly, K and $f_0(980)$, \bar{K} and $f_0(980)$, K and $a_0(980)$ and \bar{K} and $a_0(980)$ can be bound. The mechanism for binding in these cases is the S-wave kaon exchange. The nearness of one of the constituents to its decay threshold into a kaon plus a remainder, implies that the range of the kaon exchange force becomes abnormally long—significantly longer than $1/m_K$ which greatly aids the binding.

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INTRODUCTION

The particle data book abounds with hadronic resonances [1]. However, there are comparatively few states which are very close to the threshold for decay into other mesons. Recently a new near-threshold state—the narrow $D_{sJ}^*(2317)$ ($I = 0$ and possibly $J^P = 0^+$) at about 40 MeV below KD threshold—was found at BABAR [2], CLEO [3] and Belle [4].

Note, that the $D_{sJ}^*(2317)$ state can be interpreted in a number of ways: (a) as the missing triplet S-wave ($J^P = 0^+$) $|c\bar{s}\rangle$ “quarkonium state”; (b) as a “single bag” $(|c\bar{s}u\bar{u}\rangle + |c\bar{s}d\bar{d}\rangle)/\sqrt{2}$ isosinglet state; or (c) as an isosinglet “molecular” bound state $(|K^+D^0\rangle + |K^0D^+\rangle)/\sqrt{2}$ of two separate hadrons. The two hadrons in the last case can be bound—just like the deuteron—by an attractive potential due to the t-channel exchange of various light mesons. The Lagrangian has “off-diagonal” terms such as $q\bar{q}$ pair creation and annihilation and/or “bag” fissioning and rejoining interconnecting states of type (a) and (b), and (b) and (c) respectively. As a result we expect that $D_{sJ}^*(2317)$ is a superposition of all three states in (a), (b) and (c). The question is then which one dominates the state $|D_{sJ}^*(2317)\rangle$.

Regardless of how one chooses to interpret the state there is one key fact about this state which will play a major role in what follows: the state is extremely close to the KD threshold. This situation parallels a case of the pseudoscalar isosinglet and isotriplet mesons— $f_0(980)$ and $a_0(980)$ —which are very close to the $K\bar{K}$ threshold. These states can correspond to any one of the three cases above provided the quark pair $c\bar{s}$ is replaced by $s\bar{s}$ [5].

In Ref. [6] one of us argued in favor of interpretations (b) and (c). The argument in [6] was based on the fact that the mass difference of approximately 20 MeV between $D_{sJ}^*(2317)$ and the state D_0^* ($J^P = 0^+$) with a mass of about 2300 MeV (BELLE [7]) is significantly smaller than an approximate 100 MeV split between any two “strangeness analogue” $X_{\bar{s}} - X_{\bar{q}}$ ($q = u, d$) mesonic or baryonic state [8]. Likewise, the isotriplet (P -wave) $s\bar{s}$ state is 40 MeV lighter than the isotriplet S -wave $\phi(1020)$ rather than being more than 350 MeV heavier, as is the case for all other nonets. This is an argument against the “minimal” interpretation of $f_0(980)$ and $a_0(980)$ states as $s\bar{s}$ pairs. To the extent that $f_0(980)$, $a_0(980)$ and $D_{sJ}^*(2317)$ are indeed of type (b) or (c) then the following prediction can be made. A “QCD inequality” [9] implies yet another pseudoscalar $c\bar{c}$ state approximately 100 MeV below the threshold [6]. This state can be discovered via the $\eta\eta_c$ decay mode in BABAR and Belle. Ordinary $c\bar{c}$ states are accounted for and such a state would have to be interpreted as being exotic.

Of course, one can take a far more agnostic position as far as the interpretation of the $D_{sJ}^*(2317)$ or the $f_0(980)$

and $a_0(980)$. Since the three interpretations were expressed in terms of model concepts rather than QCD degrees of freedom, one can argue that even in principle there is no way to distinguish between them. However one chooses to interpret these states, we can rely on the fact that they have $J^P = 0^+$ and are only very slightly below the corresponding break-up thresholds: $40 - 50 \text{ MeV}$ below KD and $10 - 20 \text{ MeV}$ below $K\bar{K}$ thresholds respectively. This fact greatly facilitates the possibility that these mesons will be bound weakly into “molecular”-like states: $|DD_{sJ}^*(2317)\rangle$ and $|Kf_0\rangle$, $|Ka_0\rangle$, $|\bar{K}f_0\rangle$, $|\bar{K}a_0\rangle$. While any of these states would be interesting to observe, the $|D_{sJ}^*(2317)D\rangle$ is of particular interest owing to the fact that by quantum numbers alone ($S = 1, C = 2$) it is manifestly exotic.

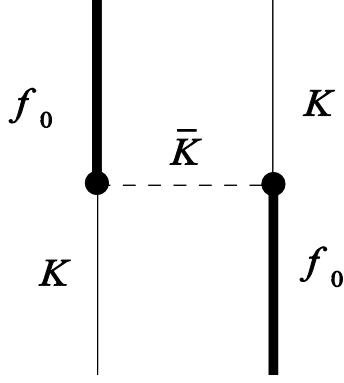


FIG. 1: $|Kf_0\rangle$ ($|\bar{K}f_0\rangle$) bound state: t -channel K exchange.

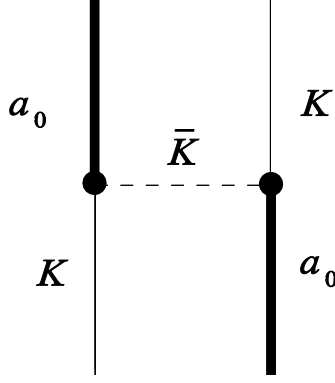


FIG. 2: $|Ka_0\rangle$ ($|\bar{K}a_0\rangle$) bound state: t -channel K exchange.

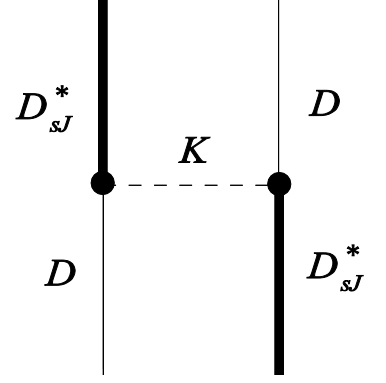


FIG. 3: $|DD_{sJ}^*(2317)\rangle$ bound state: t -channel K exchange.

THE BINDING MECHANISM

The key to our analysis is the following fact. If there exists a meson \mathcal{H}_a quite near to the break-up threshold into a K plus another meson, \mathcal{H}_b , then there will be an unnaturally long-ranged force between mesons \mathcal{H}_a and \mathcal{H}_b due to K exchange. Now suppose the K -exchange force leads to an attractive S -wave interaction. This in turn will lead to tendency toward binding. Of course, if the state formed is deeply bound, then the fact the interaction is comparatively long ranged plays no important role. Indeed, if the state is deeply bound, the K -exchange mechanism clearly does not dominate and one would need to understand the physics at the quark level. However, if the state formed is relatively loosely bound, the unnaturally long-ranged character of the potential becomes essential to the binding. As we show below, for a large range of parameters, this is precisely what occurs: the K -exchange is strong enough to bind the two mesons together but is weak enough so that the resulting bound state is dominated by ranges where the K -exchange is expected to be important. Note, that if we are in such a weakly bound region, it is legitimate to use a nonrelativistic Schrödinger equation to describe the dynamics.

We begin the analysis by deriving the interaction due to the K -exchange. Consider two “heavy” mesons \mathcal{H}_a and \mathcal{H}_b ($\mathcal{H}_a = f_0$ (a_0)—case (1)—or $D_{sJ}^*(2317)$ —case (2); $\mathcal{H}_b = K$ (\bar{K})—case (1)—or D —case (2)) of the same spin and opposite parity. The mesons need only be heavy in the sense that their masses are much larger than the ultimate binding energy between them. In this regime, the binding potential, $V(r)$ ($r = |r_a - r_b|$), can be obtained from a one-kaon exchange amplitude shown in the Feynman diagrams of Figures 1, 2, 3. They can be evaluated in the limit in which the recoil energies of the mesons \mathcal{H}_a and \mathcal{H}_b are neglected (such a limit is justified near the $\mathcal{H}_a \mathcal{H}_b$ threshold). In this limit the energy transfer carried by an off-shell K meson is $\epsilon = m_b + m_K - m_a$, where m_a , m_b , m_K are the masses of \mathcal{H}_a , \mathcal{H}_b and K mesons respectively. In the cases of interest here this energy is $\epsilon \lesssim 50 \text{ MeV}$, *i.e.* much smaller than the mass of K meson. The ϵ can be viewed as the binding energy of K and \mathcal{H}_b into \mathcal{H}_a .

Keeping the leading term in $1/m_K$ expansion of the kaon propagator and taking the Fourier transform of the amplitude one obtains an attractive Yukawa-like potential:

$$V(r) = -\frac{g_i^2}{16\pi m_a m_b} \frac{\exp[-r\sqrt{2m_K \epsilon_i}]}{r} = \alpha_i \frac{\exp[-\kappa_i r]}{r}, \quad (1)$$

where g_i ($g_1 = g_{K\bar{K}f_0(a_0)}$, $g_2 = g_{KDD_{sJ}^*(2317)}$) is (the mass dimension two) coupling constant of the S -wave Yukawa coupling $K\mathcal{H}_a\mathcal{H}_b$. The factor $4 m_a m_b$ in the denominator in Eq. (1) comes from the non-relativistic normalization of the scalar wave functions of \mathcal{H}_a and \mathcal{H}_b . Consequently, the coupling constant $\alpha_i = -g_i^2/(16\pi m_a m_b)$ is dimensionless.

The potential in Eq. (1) can be interpreted as the (asymptomatic) profile function of the field strength of the virtual K inside the H_a bound state (up to the coupling constants). It has the form of an outgoing spherical wave with a purely imaginary momentum $k_i = i \kappa_i$ with κ_i equal to $\sqrt{2 m_K \epsilon_i}$. Note that the κ_i in Eq. (1) replaces m_K in the standard Yukawa-like potential yielding much longer range potentials. This effect is huge for possible DD^* loosely state considered by Törnqvist [10]. As in our case, $\epsilon = m_D + m_\pi - m_{D^*}$ is tiny. However, in that case, the interaction is in P -wave with a derivative πDD^* coupling. As a result, the increase in the range in this case is essentially compensated by the corresponding decrease in the strength of the coupling. This is not the case for the S -wave momentum independent couplings relevant in our case.

ESTIMATED BINDING ENERGIES

The central result of this paper is that for a wide range of “reasonable” interactions between \mathcal{H}_a and \mathcal{H}_b binding results.

In the two cases considered here—(1) $|K f_0\rangle, |K a_0\rangle, |\bar{K} f_0\rangle, |\bar{K} a_0\rangle$ and (2) $|DD_{sJ}^*(2317)\rangle$ —the values of κ_i are:

$$100 \lesssim \kappa_1 \lesssim 140 \text{ MeV}, \quad 200 \lesssim \kappa_2 \lesssim 220 \text{ MeV}. \quad (2)$$

The variation in Eq. (2) is due to the differences in binding energies for various D and K charge sates.

The binding energies of the $|\mathcal{H}_a \mathcal{H}_b\rangle$ “molecules” can now be determined from the Schrödinger equation with a potential given in Eq. (1) and reduced masses $\mu_1 = 2 m_K/3 \approx 330 \text{ MeV}$ (case (1)) and $\mu_2 \approx 1030 \text{ MeV}$ (case (2)). The binding energies and the typical sizes of the ground state wave functions (given by $\sqrt{\langle r^2 \rangle}$) for a number of couplings α_i and values of κ_i (Eq. (2)) are shown in Tables I, II, III (case (1)) and Tables IV, V, VI (case (2)) [11].

In Ref. [12] the value of the coupling constant $g_{K\bar{K}f_0}^2/(4\pi)$ was determined to be 0.6 GeV^2 . The corresponding dimensionless coupling is $\alpha_{K\bar{K}f_0} = -g_{K\bar{K}f_0}^2/(16\pi m_K^2) \approx -0.6$. We also assume that the same value is applicable in the case of Ka_0 and $DD_{sJ}^*(2317)$ systems.

α	$E_B, \text{ MeV}$	$\sqrt{\langle r^2 \rangle}, \text{ fm}$
-0.4	-2.9	36.1
-0.6	-17.3	17.1
-0.8	-44.4	11.8

TABLE I: $|K f_0\rangle (|K a_0\rangle)$: the binding energies and the size of the ground sate wave function for the potential in Eq. (1) with $\kappa = 100 \text{ MeV}$.

α	$E_B, \text{ MeV}$	$\sqrt{\langle r^2 \rangle}, \text{ fm}$
-0.4	-0.3	125
-0.6	-8.3	26.7
-0.8	-28.4	15.2

TABLE II: $|K f_0\rangle (|K a_0\rangle)$: the binding energies and the size of the ground sate wave function for the potential in Eq. (1) with $\kappa = 140 \text{ MeV}$.

α	$E_B, \text{ MeV}$	$\sqrt{\langle r^2 \rangle}, \text{ fm}$
-0.2	-0.6	41.1
-0.4	-25.5	8.9
-0.6	-90.0	5.0

TABLE III: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground sate wave function for the potential in Eq. (1) with $\kappa = 200 \text{ MeV}$.

α	$E_B, \text{ MeV}$	$\sqrt{\langle r^2 \rangle}, \text{ fm}$
-0.2	-0.2	72.1
-0.4	-22.0	8.8
-0.6	-82.8	5.3

TABLE IV: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground sate wave function for the potential in Eq. (1) with $\kappa = 220 \text{ MeV}$.

For this value of the coupling constant the binding energy of $|K f_0\rangle, |K a_0\rangle, |\bar{K} f_0\rangle, |\bar{K} a_0\rangle$ systems ranges from about 8 to 20 MeV (for various values of κ), Tables I, II. In the case of $|DD_{sJ}^*(2317)\rangle$ the binding energy is about 80 – 90 MeV , Tables III, IV.

The potential given in Eq. (1) treats $f_0 (a_0)$, K , D and $D_{sJ}^*(2317)$ as though they were point-like particles. The spatial extent of the K and D mesons are approximately 0.4 and 0.3 fm . The $f_0(a_0)$ and $D_{sJ}^*(2317)$ mesons are presumably even larger (particularly if the interpretations (b) and (c) discussed in the Introduction are correct). Hence, the Yukawa potential in Eq. (1) cannot apply at distances shorter than perhaps 0.5 fm . The large width (short lifetimes), $\Gamma_{f_0/a_0} = 50 - 100 \text{ MeV}$ and $\tau \sim (1.3 - 0.7) \times 10^{-23} \text{ sec}$, makes observations of such states difficult. Roughly speaking, since $\tau < T$, with T being the time for completing one period in the bound state, $T \sim 2 m_{f_0} m_K / \kappa$, the f_0 and a_0 decay before “realizing” that they are bound. The size of $f_0 (a_0)$ is of order of 1 – 2 fm and its velocity in traversing the orbit is approximately $200 \text{ MeV} / 500 \text{ MeV} \approx 0.4$, so that $T > (2.5 - 5) \times 10^{-23} \text{ seconds}$.

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-8.71	13.7
-0.6	-27.9	8.5
-0.8	-53.3	6.4

TABLE V: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 200 MeV$, $R = 0.5 fm$.

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-14.8	10.5
-0.6	-46.9	6.5
-0.8	-91.0	5.4

TABLE VII: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 200 MeV$, $R = 0.3 fm$.

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-6.5	17.6
-0.6	-23.3	8.9
-0.8	-46.3	7.0

TABLE VI: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 220 MeV$, $R = 0.5 fm$.

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-4.9	17.5
-0.6	-16.7	10.2
-0.8	-32.3	7.9

TABLE VIII: $|D D_{sJ}^*(2317)\rangle$: the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 200 MeV$, $R = 0.7 fm$.

In the case of the $|DD_{sJ}^*(2317)\rangle$ “molecule” the spatial extent of the stable D and a very long lived $D_{sJ}^*(2317)$ [13] can be expected to reduce the binding energies. In this case the range of the Yukawa potential, Eq. (1), is approximately $1 fm$. The most conservative approach to the unknown short range physics is to cut off the Yukawa potential at distances shorter than, say, $R = 0.5 fm$ and assume that $V(r) = V(R)$ (the value of the potential in Eq. (1) at $r = R$) for $r < R$. The binding energies and the corresponding sizes of the wave functions are shown in Tables V, VI, VII, VIII, IX, X. As can be expected, there is a reduction in the binding energies. The $|DD_{sJ}^*(2317)\rangle$ system is still bound by about $20 - 30 MeV$. The size of the bound state—of order of $8 fm$ is dominated by the tail of the K -exchange potential, Eq. (1). We note in passing that in both cases the systems become unbound if the potential is taken to be zero at $r < R$ (a radical assumption).

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-2.4	40.7
-0.6	-14.3	20.4
-0.8	-35.2	13.8

TABLE IX: $|K f_0\rangle$ ($|K a_0\rangle$): the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 100 MeV$, $R = 0.3 fm$.

α	E_B, MeV	$\sqrt{\langle r^2 \rangle}, fm$
-0.4	-1.9	48.3
-0.6	-11.4	20.6
-0.8	-27.6	14.2

TABLE X: $|K f_0\rangle$ ($|K a_0\rangle$): the binding energies and the size of the ground state wave function for the potential in Eq. (1) for $r > R$ and $V(r) = V(R)$ for $r < R$; $\kappa = 100 MeV$, $R = 0.5 fm$.

Note that as the composite states overlap the strong short range hyperfine interactions come into play since we have largely different quarks in $D_{sJ}^*(2317)$ and D even for the case (b) above with $D_{sJ}^*(2317)$ viewed as a four-quark construct. The tendency to form these new loosely bound states would imply that at shorter distances we have even stronger attraction than the extrapolation of the relatively smooth Yukawa potential to short distances and the results without any cutoff and a fortiori those in the case (2) may be relevant!

It is interesting to note the drastic consequence of an even small attractive scattering length—with no bound state in KK (rather than $K\bar{K}$ channel). Arbitrary (sufficiently large) number of K^0 in a common S -wave state would then attract forming a condensate carrying macroscopic strangeness *ala* Lee and Yang or Coleman’s Q -balls [14]. The longest range interaction between two kaons (and in fact any two mesons!) due to the two pion exchange—specifically the S -wave projection thereof in the t -channel is like a σ ($J^{PC} = 0^{++}$) or a scalar graviton exchange which is always attractive. The same also holds for KN interactions. However, the scattering length in the Born approximation appropriate here is given by $\int dr r^2 V(r)$ and the long range attraction is overcome (surely for KN from scattering data analysis and most likely for KK) by the strong short range repulsion so that the condensates may not exist.

It is amusing to note in passing the (admittedly weak) connection between the $|D D_{sJ}^*(2317)\rangle$ bound state and the “Efimov effect” [15]. The latter (which inspired us to look at the present problem) would arise for a zero energy

$|KD\rangle$ S -wave bound state and infinite scattering length. This in turn leads to an infinite series of three body $|KDD\rangle$ bound states. The ratios the binding energies and the sizes of the Efimov states scale as $E_B(n+1)/E_B(n) = e^{-2\pi} \sim 0.0016$, $\langle r \rangle (n+1)/\langle r \rangle (n) = e^\pi \sim 25$ (see also [16, 17, 18]). Clearly in the present case where the range of the actual potential is only about 1 fm this idealized case and the very extended— $\langle r \rangle (n) \sim 25^n$ (or contracted)—states in the above series are irrelevant. Note, the the Yukawa potential, Eq. (1) goes to $1/r$ in the limit as ϵ goes to infinity and κ goes to zero rather rather than $1/r^2$ as in the Efimov effect. The reason is that the Efimov effect requires exact diagonalization of the degenerate perturbation transcending the perturbative one-meson exchange [17, 18].

EXPERIMENTAL SIGNATURES

Assuming that the $|DD_{sJ}^*(2317)\rangle$ “molecular” state exists how can it be produced and detected? Since the production requires two pairs of $c\bar{c}$ quarks the discussion in [19] is relevant, providing an upper bound on the expected rate of the new loosely bound extended $|DD_{sJ}^*(2317)\rangle$ state which we term \mathcal{M} for *molecular*.

The bound state should manifest as a narrow peak in the mass distribution of associate D and $D_{sJ}^*(2317)$ decays. However, the limited experimental resolution at BABAR and Fermi Lab experiments limits the extent that we can utilize this. The different binding of D^+ and D^0 and different life-time $\tau_{D^0} \sim 0.5 \tau_{D^+}$ may lead to some extra signatures in specific charge dependence of the width and even in the binding energies.

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